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A new mechanism for baryogenesis living through electroweak era

Hyung Do Kim[(a)*], Jihn E. Kim[(a)†] and Takuya Morozumi[(b)‡]
[(a)] *School of Physics, Seoul National University, Seoul 151-747, Korea*
[(b)] *Graduate School of Science, Hiroshima University, Higashi Hiroshima–739–8526, Japan*

Abstract

We present a new mechanism for baryogenesis by introducing a heavy vector-like $\mathrm{SU}(2)_W$ singlet quark(s) with $Q_{\rm em} = \frac{2}{3}$ quark U or $Q_{\rm em} = -\frac{1}{3}$ quark D. The lifetime of the heavy quark is assumed to be in the range 2×10^{-11} s $< \tau <$ 1 s. Being $\mathrm{SU}(2)_W$ singlet, it survives the electroweak phase transition era. It mixes with $\mathrm{SU}(2)_W$ doublet quarks with tiny mixing angles to satisfy the FCNC constraints, where a simple Z_2 symmetry is suggested for realizing this scheme. The heavy quark asymmetry is generated in analogy with the old GUT scenario.

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*hdkim@phya.snu.ac.kr

†jekim@phyp.snu.ac.kr

‡morozumi@hiroshima-u.ac.jp

The baryon asymmetry of universe is one of the most important cosmological observations needed to be explained by a particle physics model(s). Starting from a baryon symmetric universe, Sakharov proposed three conditions for generating a baryon asymmetry in the universe from fundamental interactions applicable in cosmology [1]: the existence of baryon number violating interactions which accompany C and CP violation, and their working in a non-equilibrium state in the cosmos. With the advent of grand unified theories(GUTs), some GUTs applied in the evolution of the universe satisfied these three conditions [2]. A popular scenario was to have superheavy colored scalars($\mathbf{5}_H$s in SU(5) for example) decaying at least to two channels, say qq and ql, both of which violate the baryon number. Until mid-1980s, this GUT scenario seemed to be the theory for the baryon asymmetry in the universe.

However, this GUT scenario underwent a nontrivial modification after considering high temperature effects in nonabelian gauge theories. The standard model(SM) has a non-perturbative sphaleron solutions [3] whose effect on tunneling is supposed to be extremely small, $\sim e^{-8\pi^2/g_W^2}$. But at high temperature the sphaleron barrier for tunneling is overcome and it was pointed out that the tunneling through sphaleron can be important in cosmology [4]. Since the sphaleron violates the baryon number, it was argued that the baryon asymmetry produced in the earlier epoch is erased during the epoch of electroweak phase transition. In the GUT scheme, some models can keep an asymmetry [5]. One example is to use the $B-L$ conservation of SO(10). Using the $B-L$ conservation during the electroweak phase transition, the leptogenesis generating the B number from a nonvanishing L number was proposed [6]. The most recent complete phenomenological analysis on the leptogenesis constrains the lightest singlet Majorana neutrino mass $M_1 > 10^9$ GeV [7], probably contradicting the upper bound of the reheating temperature after inflation in supergravity models [8]. Therefore, theoretically and also for a practical implementation of the baryogenesis, it is desirable to invent any new mechanism for baryogenesis.

In this paper, we devise a mechanism for evading the sphaleron ΔB erasing scenario. The mechanism employs a vector-like $\mathrm{SU}(2)_W$ singlet heavy quark Q, i.e. Q_L and Q_R, which can be $Q_{\rm em} = \frac{2}{3}$ quark U [9] or $Q_{\rm em} = -\frac{1}{3}$ quark D [10] so that it decays to ordinary quarks through electroweak interactions. The $\mathrm{SU}(2)_W$ singlet vector-like quark(s) necessarily introduces flavor changing neutral currents [9,11] which is harmful if the heavy quark mass is below the weak scale. If the heavy quark is sufficiently heavy and survives through the epoch of the electroweak phase transition, then the heavy quark asymmetry generated by the Sakharov conditions would remain unwashed by the sphaleron processes. This is because

the 't Hooft interaction for an $\mathrm{SU}(2)_W$ sphaleron does not involve $\mathrm{SU}(2)_W$ singlets. Being $\mathrm{SU}(2)_W$ singlets, a heavy quark asymmetry would not be erased. The asymmetry in heavy quarks may be erased by the mixing with $\mathrm{SU}(2)_W$ doublet quarks. But the mixing angle ϵ in our scheme is so small that the washing out of a heavy quark asymmetry by the sphaleron is completely negligible. We will discuss more on this later. Our proposal depends on whether one can construct a theoretically reasonable model with $\mathrm{SU}(2)_W$ singlet heavy quarks with the following properties:

- The Sakharov conditions for creating $\Delta Q \sim \Delta B$ are available. For $\Delta B \neq 0$ processes, GUTs are used. Also the CP violation at the GUT scale is present, and the heavy colored scalar particles are present for them to decay to colored quarks, anti-quarks and/or leptons.

- The lifetime of Q is sufficiently long, $\tau_Q > O(10^{-11}\ \mathrm{s})$ [13], so that the Q asymmetry survives the chaotic electroweak phase transition.

- The heavy quark model must not introduce too large flavor changing neutral currents.

For definiteness, let us introduce $Q_{\rm em} = -\frac{1}{3}$ quarks Ds.

The first point is easily implementable in GUTs with $\mathrm{SU}(2)_W$ singlet quarks such as in an E_6 GUT. In the E_6 GUT, one family is embedded in $\mathbf{27}_F$ which can contain 15 chiral fields of the SM, a vector-like lepton doublets $\{L_1(Y=-\frac{1}{2}), L_2(Y=\frac{1}{2})\}$, one vector-like heavy quark $D(Y=-\frac{1}{3})$ and $D^c(Y=\frac{1}{3})$, and two heavy neutrinos. For three families, we introduce three $\mathbf{27}_F$s. For the Higgs mechanism, we introduce a scalar $\mathbf{27}_H$ which contains three Higgs doublets. An adjoint representation $\mathbf{78}_H$ is needed for breaking E_6 down to the SM. Certainly, there exists a gauge hierarchy problem of how we remove most scalars of $\mathbf{27}_H$ at the GUT scale, which is not addressed here. In this setup, we have all the ingredients for the GUT baryogenesis [2]. Below, however, we will not restrict to the E_6 GUT, but proceed to discuss in the SM framework with a heavy $Q_{\rm em} = \frac{1}{3}$ colored scalar X_i (anti-fundamental of $SU(3)_C$) to generate the heavy quark D number asymmetry ΔQ with $i = 1, 2$. Relevant interactions are

$$g_{Di} X_i u^c D^c + g_{ei} X_i^* u^c e^c + \text{h.c.}. \tag{1}$$

All the fermions are written in terms of left-handed Weyl spinor. (ψ^c is the charge conjugation of the right-handed spinor.) The diagrams responsible for the ΔQ asymmetry in the decay of heavy colored scalar fields are shown in Fig. 1.

FIGURES

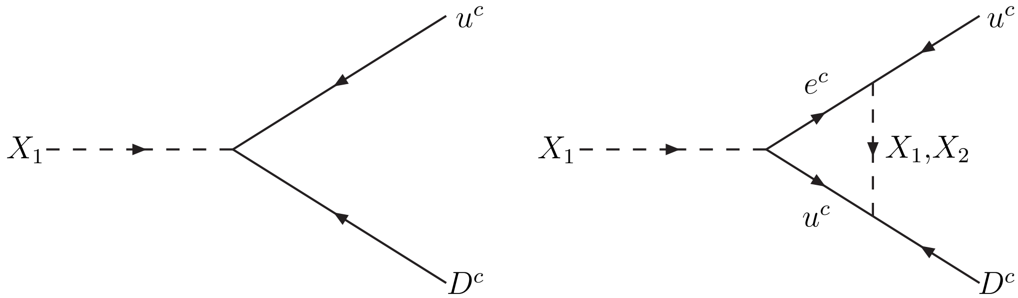

FIG. 1. The interference between tree and one-loop diagrams is needed for a nonzero ΔQ generation.

The decay of X_i (X_i^*) gives a positive (negative) ΔQ number. We can consider all possible baryon number generations, but those light-quark numbers are washed out during the electroweak phase transition through the sphaleron process and only ΔQ number is the meaningful conserved quantity.

If we had just a single colored scalar X ($i = 1$), the first and the second diagrams in Fig. 1 are proportional to g_D and $g_e^* g_e g_D$, respectively, and the crossing term $g_e^* g_e g_D^* g_D$ is real. In this case we should consider higher loop corrections which are highly suppressed [12]. The leading interference term can be complex if we have two colored scalars X_i with $i = 1, 2$. The phase from this leading interference term is proportional to

$$\arg\left(g_{D1}^* g_{D2} g_{e1}^* g_{e2}\right). \tag{2}$$

If we allow arbitrary phases in the Yukawa couplings, the relative phase of g_{D1} and g_{D2} can be canceled only by the relative phase redefinition of X_2 compared to X_1. It is also true for g_{e1} and g_{e2}. Therefore, if $\arg(g_{D1}/g_{D2})$ is different from $\arg(g_{e1}/g_{e2})$, one of the phases in the Yukawa couplings cannot be rotated away. This proves that the phase appearing in the interference (2) is physical (i.e. *un-removable*).

In principle, we can consider decays of both X_1 and X_2. However, we will assume that X_2 is heavier than X_1 and the decay of X_1 dominates.

ΔQ number generated from the above decay is, including the contribution from the wave function correction diagram [14],

$$\frac{n_D}{s} \simeq \frac{\kappa}{4\pi g_*} \frac{\mathrm{Im}(g_{D1}^* g_{D2} g_{e1}^* g_{e2})}{(g_{D1}^* g_{D1} + g_{e1}^* g_{e1})} \left[f_v(x) + f_s(x)\right], \tag{3}$$

where κ is the washout factor, $f_v(x) = 1 - x\log(1 + \frac{1}{x})$, $f_s(x) = \frac{1}{x-1}$, $x = \frac{m_{X_2}^2}{m_{X_1}^2}$ is the mass ratio of two heavy scalar fields and $g_* = 106.75$ is the number of effective degrees of freedom

in the SM. Within the range of parameters we discuss, it is possible to obtain the observed value $n_B/s = (8.7 \pm 0.4) \times 10^{-11}$ in our scheme. Generation of ΔQ number after inflation will be discussed later.

For a very long lifetime τ_D, the mixing angle between the ordinary quarks and heavy quarks D must be sufficiently small. The estimation of the mixing angle dependence on the quark masses can be considered in the following way. Introducing a discrete symmetry, e.g. by giving different discrete quantum numbers to b_R and D_R, one can consider the following 2×2 simplified mass matrix for an ordinary quark b and a heavy singlet quark D,

$$A = \begin{pmatrix} m & J \\ 0 & M \end{pmatrix} \tag{4}$$

where J parametrizes the D_R coupling to b_L. The eigenvalues of $AA^\dagger$ are found to be

$$\begin{pmatrix} |m_b|^2 \\ |m_D|^2 \end{pmatrix} = \frac{1}{2}\left(|M|^2 + |m|^2 + |J|^2\right) \mp \frac{1}{2}\sqrt{[(|M| + |m|)^2 + |J|^2][(|M| - |m|)^2 + |J|^2]} \tag{5}$$

which, in the limit $|M|^2 \gg |m|^2, |mJ|$, approximate to

$$|m_b| \simeq |m|, \quad |m_D| \simeq |M| \tag{6}$$

The Hermitian matrix $AA^\dagger$ is diagonalized by a unitary matrix. Taking vanishing phases, in the super-large M limit we obtain the eigenstates,

$$|b\rangle \simeq \begin{pmatrix} 1 \\ -\frac{J}{M} \end{pmatrix}, \quad |D\rangle \simeq \begin{pmatrix} \frac{J}{M} \\ 1 \end{pmatrix}. \tag{7}$$

From Eq. (7), one can make the mixing sufficiently small by taking $J/M \to 0$.

For generalizing to three ordinary quarks d_i $(i = 1, 2, 3)$ and n heavy quarks D_J $(J = 1, \cdots, n)$, one can consider the following $(3 + n) \times (3 + n)$ mass matrix $\mathbf{M}$ to

$$\mathbf{M} = \begin{pmatrix} M_d & J \\ J' & M_D \end{pmatrix}$$

where M_d is a 3×3 mass matrix for ordinary $Q_{\rm em} = -\frac{1}{3}$ down-type d quarks , J is a $3 \times n$ matrix, J' is an $n \times 3$ matrix, and M_D is an $n \times n$ mass matrix for the heavy $Q_{\rm em} = -\frac{1}{3}$ quarks Ds. We can redefine the right-handed fields only to make J' vanish. If J'/M_D is small and $M_d \sim J$, there is no correction to the matrix. Thus without loss of generality, we can consider the following mass matrix

$$\mathbf{M} = \begin{pmatrix} M_d & J \\ \mathbf{0} & M_D \end{pmatrix} \tag{8}$$

where $n \times 3$ elements of matrix $\mathbf{0}$ are zeros. If the elements of M_d are O(m), the elements of J are O(J), and the elements of M_D are O(M), the mixing angles between d and D quarks are of order $\epsilon \sim J/M$. Thus, if ϵ is sufficiently small then the lifetime(s) of D quark(s) can be made long.

For an order of magnitude estimation for the lifetime of the lightest D, let us use the $b - D$ system with negligible couplings to the other light quarks, viz. Eq. (4). The flavor changing neutral current coupling of Z_μ is of order ϵ also. Thus, we can estimate the decay width of D from $D \to tW, bZ, bH^0$ as

$$\Gamma_D = \frac{\sqrt{2}G_F}{8\pi} J^2 m_D. \tag{9}$$

where we assumed $m_D \gg m_t$. J is parametrized as

$$|J| \equiv f m_b \tag{10}$$

with f a small number. Here, $\epsilon \equiv f m_b/m_D$. The lifetime of D should be made longer than 2×10^{-11} s for D to pass through the electroweak phase transition era. However, it should not be too long, say $\tau_D < 1$ s, not to disrupt the standard nucleosynthesis. Therefore, we require the following *cosmological lifetime window* for the lightest D,

$$2 \times 10^{-11} \text{ sec} \le \tau_D \le 1 \text{ sec} \tag{11}$$

The lifetime window given above can be translated into the constraint equation on $|\epsilon|$

$$\frac{1}{(10^6\ m_D(\text{GeV}))^{3/2}} \le |\epsilon| \le \frac{1}{(2.7 \times 10^2\ m_D(\text{GeV}))^{3/2}}. \tag{12}$$

For a large m_D, on the other hand, the four fermion interaction is the dominant one. Then, the condition (D decay rate) $\le H$ at $T \sim M_W$ is $m_D \le 10^5 (M_{X_1}/10^{10}\text{ GeV})^{4/5}$ GeV.

To estimate how much of ΔD is washed out by the sphaleron processes, we first estimate the mixing of mass eigenstates b_L and D_L in the weak eigenstate b_L^0: $b_L^0 \simeq b_L + \epsilon D_L$. The oscillation period is $\sim 1/m_D$. The probability to find SU(2)$_W$ doublet b_L^0 from D_L is $\sim |\epsilon|^2$ in the time interval $1/m_D$, or the rate to go to b_L^0 is $m_D|\epsilon|^2$. Since the period of the electroweak phase transition lasts for $1/H \sim M_P/M_W^2$, the amount of ΔQ washed out during the electroweak phase transition is $(m_D M_P/M_W^2)|\epsilon|^2$. This condition gives a rough bound, $|\epsilon| \le 10^{-8}$ for $m_D \simeq (100 - 1000)$ GeV. For $m_D > 1$ TeV, Eq. (12) gives a stronger bound.

The model presented above is constrained by the FCNC and proton decay experiments. The mass of D is constrained by the bounds on FCNC.

The mixing of D quark with b quark may change the flavor diagonal coupling $Z \to b\bar{b}$ from the standard model one

$$z_{bb} = 1 - |\epsilon|^2. \tag{13}$$

The experimental bound on $z_{bb} = 0.996 \pm 0.005$ leads to $|\epsilon|^2 \leq 0.009$. [15] If D quark mixes with b and s with the strength J_b and J_s, rare B decay $B \to Xsl^+l^-$ occurs in the tree level. The bound obtained from the analysis [16]

$$|z_{sb}| = \frac{J_b J_s}{m_D^2} < 1.4 \times 10^{-3}. \tag{14}$$

Finally, we consider FCNC constrains obtained from Kaon system. We require the FCNC contribution is smaller than the standard model. From the K_L and K_S mass difference mass difference,

$$|z_{sd}| \leq \left(\frac{G_F m_c^2 \lambda^2}{2\sqrt{2}\pi^2}\right)^{1/2}. \tag{15}$$

From the $K^+ \to \pi^+ \nu\bar{\nu}$

$$\frac{Br(K^+ \to \pi^+\nu\bar{\nu})|_{FCNC}}{Br(K^+ \to \pi^0 e^+ \nu)} = \frac{3}{2}\frac{|z_{sd}|^2}{\lambda^2} \leq 2 \times 10^{-9} \tag{16}$$

This leads to

$$|z_{sd}| \leq 7.3 \times 10^{-6}. \tag{17}$$

If we neglect the flavor dependence of J, the upper bounds on ϵ obtained from $Z \to b\bar{b}$, $B \to Xsl^+l^-$, Δm_K and $K^+ \to \pi^+\nu\bar{\nu}$ are 0.037, 0.095, 0.02 and 7.3×10^{-6}, respectively. The tightest constraint from $K^+ \to \pi^+\nu\bar{\nu}$, gives $m_D \geq 6.6 \times 10^6 f$ GeV from $|\epsilon| \simeq f m_b/m_D$. Then, Eq. (12) gives

$$\frac{1}{4.8 \times 10^9 \sqrt{m_D(\mathrm{GeV})}} < |f| < \frac{1}{2.1 \times 10^4 \sqrt{m_D(\mathrm{GeV})}} \tag{18}$$

which can be satisfied by some range of small couplings. The J term is supposed to arise from breaking a Z_2 symmetry, and its smallness can be implemented naturally.

The highly suppressed J can be obtained by introducing a discrete symmetry. Consider a Z_2 symmetry under which the quarks have the following charges (D parity),

$$Z_2: \; b_{L,R} \to b_{L,R}, \; D_{L,R} \to -D_{L,R}. \tag{19}$$

This implies that the SU$(2)_W$ doublet q_L housing b_L has the same $Z_2 = +1$ eigenvalue. The SU$(2)_W$ singlet D can have a bare mass or obtain a mass by a large VEV of $Z_2 = +1$ singlet scalar S. All the interactions are consistent with D parity if L, e^c and X are D parity odd and all the other fields are even. We introduce two Higgs doublets, ϕ and φ,

$$Z_2 : \ \phi \to \phi, \ \varphi \to -\varphi. \tag{20}$$

If this Z_2 symmetry were exact, we obtain $J = 0$. Here we introduce a small amount of Z_2 breaking by a soft term m_δ^2,

$$\begin{aligned} V(\phi, \varphi) = (m_\delta^2 \varphi^\dagger \phi + h.c.) - \mu^2 \phi^\dagger \phi + M_\varphi^2 \varphi^\dagger \varphi + \lambda_1 (\varphi^\dagger \varphi)^2 + \lambda_2 (\phi^\dagger \phi)^2 \\ +\lambda_3 \varphi^\dagger \varphi \phi^\dagger \phi + (\lambda_4 \varphi^\dagger \phi \phi^\dagger \varphi + \lambda_5 \varphi^\dagger \phi \varphi^\dagger \phi + h.c.) \end{aligned} \tag{21}$$

where we assume $M_\varphi^2 \gg \mu^2 \gg |m_\delta^2| > 0$. Then, ϕ develops a VEV v for the electroweak symmetry breaking, but φ does not if $m_\delta^2 = 0$. For a nonzero m_δ^2, φ develops a tiny VEV, $\langle \varphi \rangle \simeq v m_\delta^2 / M_\varphi^2$. Thus, the coupling $f_{\rm off} \bar{q}_L \varphi D_R$ gives $J \sim f_{\rm off} v m_\delta^2 / M_\varphi^2 \sim (\sqrt{2} f_{\rm off}/f_b) m_b m_\delta^2 / M_\varphi^2$. Then, our small parameter f is $f = (\sqrt{2} f_{\rm off}/f_b) m_\delta^2 / M_\varphi^2$. A small soft mass can lead to a very small f even for O(1) value of $f_{\rm off}$. One loop diagram can generate J which must be proportional to m_δ^2; hence it is subdominant. By the choice of the soft Z_2 breaking term, J can be made to fall in the region for the needed lifetime window (11) for D.

Assuming no approximate discrete symmetry, proton decay can proceed via the X particle exchange $d^c u^c \to X^* \to ue$. Thus, the mass of X particle should be in the GUT scale with a small Yukawa couplings to the first family members as studied in GUT proton decay. This is the standard colored Higgs mediated proton decay. But the mass of D is not restricted by proton decay.

But in the inflationary scenario, the X particle should be light enough($< 10^{13}$ GeV) so that enough Xs are present after inflation. In supersymmetric models, one may need a stronger constraint, $M_X < 10^{9-10}$ GeV, from the gravitino problem [8]. To implement this constraint, one can use the above softly broken Z_2 symmetry with X carrying $Z_2 = -1$ parity. Thus, the proton decay operator $ud \to X \to u^c e^+$ has an additional suppression factor ξ^2 which is nonvanishing only if the Z_2 is broken. For $\tau_p > 10^{33}$ years, we obtain $\xi < 0.7 \times 10^{-6}$ for $M_X = 10^{10}$ GeV. Preheating scenario [17] can generate sizable baryon asymmetry even for heavy X (heavier than inflaton). It can happen if $\Gamma_X \le 10^{-3} M_X$ which can be satisfied for $g_{1(e,D)} \sim 10^{-1}$. Therefore, we can generate sizable baryon asymmetry from X decay after inflation.

In conclusion, we devised a new mechanism for baryogenesis by introducing $SU(2)_W$ singlet quark(s). The conditions for the current mechanism to work are to generate ΔQ at the GUT scale and the lifetime bound must fall in the region (11).

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